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Existence of nuclei with unusual neutron excess?

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Abstract

A realistic model is suggested based on the quasiparticle Lagrange version of the self-consistent Finite Fermi Systems theory supplemented with the microscopically calculated surface parameters of the Landau–Migdal interaction amplitude. The latter are expressed in terms of the off-shell T-matrix of free NN-scattering and show a strong dependence on the chemical potential of a nucleus under consideration in the drip line vicinity. This effect could result in shifting the neutron drip line position to very large values of the neutron excess.

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Up to now, all predictions on the location of the nuclear drip line for heavy nuclei are based on phenomenological approaches. All of them operate with parameters adjusted to the properties of stable nuclei. This point seems to be questionable for nuclei far from the β-stability valley. In view of commonly recognized importance of pairing for the drip line, the main efforts in this field were focused to study the superfluidity effects at small values of the chemical potential μ [1–3]. Such problems were investigated as a correct account of continuum, comparison of exact Bogolyubov solutions versus those within the BCS approximation, the surface pairing versus the volume one, and so on. In this Letter we concentrate on the examination of the neutron average potential well at small values of μ_n. We give arguments in favour of significant variation, in the vicinity of the drip line, of some parameters of the effective NN-interaction which generates this potential well. It could result in shifting the position of the drip line into the region of large (or, maybe, very large) values of the neutron excess.

The reason for such an effect can be readily shown in terms of the simplest version of the self-consistent (SC) Finite Fermi Systems (FFS) theory [4, 5], which is based on the simplified version of the self-consistency relation of Ref. [6]:

$$\frac{\partial U}{\partial \mathbf{r}} = \int F(\mathbf{r}, \mathbf{r}') \frac{\partial \rho}{\partial \mathbf{r}'} d\mathbf{r}', \tag{1}$$

where F is the Landau–Migdal (LM) amplitude. Obvious isotopic indices in Eq. (1) are omitted. Let us for a while limit ourselves to the zero-range components of F, which in the standard FFS theory

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notation has the form:

$$F(\mathbf{r},\mathbf{r}') = C_0\big[f_0(\mathbf{r}) + f_0'(\mathbf{r})\tau_1\tau_2 + \big(g_0(\mathbf{r}) + g_0'(\mathbf{r})\tau_1\tau_2\big)\sigma_1\sigma_2\big]\delta(\mathbf{r}-\mathbf{r}'), \tag{2}$$

where the normalization factor $C_0 = (dn/d\varepsilon_{\mathrm{F}})^{-1}$ is the inverse density of states at the Fermi surface.

In the FFS theory, a strong dependence of the scalar-isoscalar amplitude f_0 on the observation point $\mathbf{r}$ was introduced. In fact, in Ref. [4] a simple interpolation form of such a dependence was suggested:

$$f_0(\mathbf{r}) = f^{\mathrm{ex}} + \big(f^{\mathrm{in}} - f^{\mathrm{ex}}\big)\frac{\rho_+(\mathbf{r})}{\rho_0}. \tag{3}$$

Here $\rho_+(\mathbf{r}) = \rho_n(\mathbf{r}) + \rho_p(\mathbf{r})$ is the nuclear density in the point $\mathbf{r}$, while $\rho_0 = \rho_+(r=0)$. The subscript "0" for the zero Landau harmonics is for brevity omitted in the r.h.s. of Eq. (3) and below.

It is worth to mention that the density dependence of the phenomenological Skyrme force [7] agrees with the ansatz of Eq. (3). There exist also alternative versions of the interpolation formula for $f_0(\mathbf{r})$ by means of the function $(\rho/\rho_0)^{\alpha}$, $\alpha \neq 1$, in the r.h.s. of Eq. (3) or a more complicated density dependence [3, 8–10]. But all of them are characterized by a strong difference between the dimensionless parameters f^{ex} and f^{in}: $f^{\mathrm{ex}} \simeq -3$, whereas f^{in} is close to zero.[1]

It should be noted that the density dependent scalar–isoscalar amplitude

$$f_0'(\mathbf{r}) = f'^{\mathrm{ex}} + \big(f'^{\mathrm{in}} - f'^{\mathrm{ex}}\big)\frac{\rho_+(\mathbf{r})}{\rho_0}, \tag{4}$$

was used in a new version of the energy functional method by S. Fayans et al. [3,10] based on a detailed analysis of long isotopic chains. The difference between the parameters f'^{in}, f'^{ex} is also significant, though not so dramatic as in the isoscalar case. In principle, relations similar to Eq. (3) can be written also for the spin-dependent terms of Eq. (2), but up to now no evidence of a noticeable difference between the internal and external values of these amplitudes was found. Therefore, the equalities $g^{\mathrm{in}} = g^{\mathrm{ex}}$, $g'^{\mathrm{in}} = g'^{\mathrm{ex}}$ were imposed in the FFS theory.

In Ref. [11] it was found that the external values of the LM amplitudes can be calculated in terms of the off-shell T-matrix of free NN-scattering taken at a negative energy $E = 2\mu$, where μ is the chemical potential of the nucleus under consideration. Stable nuclei with $\mu_n = \mu_p \simeq -8$ MeV were considered, and a reasonable agreement with the phenomenological values for the surface parameters of the LM amplitude was obtained.[2] It confirmed the relevance of the asymptotic relation $F \to T(2\mu)$ for the description of the properties of stable nuclei.

Let us mention that, although the interpolation ansatz of Eq. (3) looks very alike to the LDA prescription, it cannot be obtained microscopically within the LDA. Indeed, there is a domain of density values in the surface region for which infinite nuclear matter is unstable. Within the Brueckner theory, a more consistent way to obtain the LM amplitude in a form similar to Eq. (3) should involve, first, a direct calculation of the G-matrix for the non-uniform system in the coordinate representation, and, second, application of some reasonable localization recipe to the non-local G-matrix (see, e.g., Refs. [13,14]). For the spin independent amplitudes under consideration the explicit form of these relations is as follows:

$$f_0^{\mathrm{ex}} = \frac{3}{16}\big[t_0(E=2\mu) + t_1(E=2\mu)\big], \tag{5}$$

$$f_0'^{\mathrm{ex}} = \frac{1}{16}\big[t_0(E=2\mu) - 3t_1(E=2\mu)\big], \tag{6}$$

where t_0, t_1 are the dimensionless values of the off-shell T-matrix with the spin value $S=0$ and $S=1$, respectively, taken at the zero value of all nucleon momenta. In stable nuclei the isospin symmetry works well and one has:

$$f_{nn}^{\mathrm{ex}} = f_{pp}^{\mathrm{ex}} = f_0^{\mathrm{ex}} + f_0'^{\mathrm{ex}} = \frac{1}{4}t_0(E=2\mu), \tag{7}$$

$$f_{np}^{\mathrm{ex}} = f_0^{\mathrm{ex}} - f_0'^{\mathrm{ex}} = \frac{1}{8}\big[t_0(E=2\mu) + 3t_1(E=2\mu)\big]. \tag{8}$$

In this Letter these relations are applied to nuclei close to the drip line in which the neutron and proton chemical potentials deviate significantly from each other. As far as both amplitudes t_0 and t_1 at small energy E depend on E rather sharply, the isospin

[1] The inequality $f^{\mathrm{in}} > -0.5$ should be fulfilled to avoid the Pomeranchuck instability [4].

[2] They depend a little on the type of the density dependence. The best agreement was achieved with the spin independent parameters of Ref. [3] and the spin dependent parameters of Ref. [12].

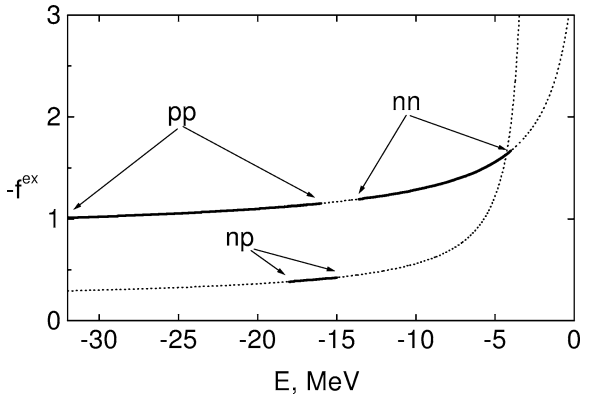

Fig. 1. External LM amplitudes $f_{nn}^{\rm ex}(E)$, $f_{pp}^{\rm ex}(E)$, and $f_{np}^{\rm ex}(E)$ (the latter is divided by 10) taken with negative sign. The solid lines show these amplitudes within the energy limits which are typical for tin isotopes.

symmetry is destroyed. As a result, one deals with the situation when $f_{nn}^{\rm ex} \neq f_{pp}^{\rm ex}$. Both of these amplitudes can be found from Eq. (7), but at $E = 2\mu_n$ for the first one and at $E = 2\mu_p$, for the second one. The amplitude $f_{np}^{\rm ex}$ can be obtained from Eq. (8) at $E = \mu_n + \mu_p$. In the neutron drip line vicinity, only the neutron chemical potential is close to zero, whereas, on the contrary, $|\mu_p|$ increases with approaching the boundary. Therefore, only the amplitude $f_{nn}^{\rm ex}$ changes significantly (see Fig. 1).

The SC relation, Eq. (1), with the amplitude F given by Eqs. (2)–(4) can be readily integrated over for a spherical nucleus yielding the following relation for the neutron potential:

$$U_n(r) = C_0\Bigg[f_{nn}^{\rm ex}\rho_n(r) + \big(f_{nn}^{\rm in} - f_{nn}^{\rm ex}\big)\frac{\rho_n(r)}{2\rho_0}\big(\rho_+(r) + \rho_p(r)\big) + f_{np}^{\rm ex}\rho_p(r) + \big(f_{np}^{\rm in} - f_{np}^{\rm ex}\big)\frac{\rho_p(r)}{2\rho_0}\big(\rho_+(r) + \rho_p(r)\big)\Bigg]. \tag{9}$$

A similar relation is obtained for the proton potential $U_p(r)$ replacing the indices "n" and "p".

To estimate the effects of the energy (or μ-) dependence of the interaction amplitude $f_{nn}^{\rm ex}$, let us find the value of $U_n(0)$. To simplify the expression, we use the approximations $\rho_n(0) = (N/A)\rho_0$ and $\rho_p(0) = (Z/A)\rho_0$. Then we get:

$$U_n(0) = \frac{1}{2}C_0\rho_0\Bigg[f_{nn}^{\rm ex}\frac{N^2}{A^2} + f_{nn}^{\rm in}\left(1 - \frac{Z^2}{A^2}\right) + f_{np}^{\rm ex}\frac{Z^2}{A^2} + f_{np}^{\rm in}\left(1 - \frac{N^2}{A^2}\right)\Bigg]. \tag{10}$$

Let us add a small number of neutrons to a heavy nucleus. Then, neglecting for a while the pairing effects, we obtain an approximate relation

$$\delta\mu_n = \delta U_n(0) \tag{11}$$

for the change of the neutron chemical potential.[3] The variation of the expression (10) is a sum of two terms,

$$\delta U_n(0) = \delta U_1 + \delta U_2, \tag{12}$$

having different origin. The first term results from variation of N and A while the second one,

$$\delta U_2 = \frac{1}{2}C_0\rho_0\frac{N^2}{A^2}\delta f_{nn}^{\rm ex}, \tag{13}$$

is due to the μ-dependence of the amplitude $f_{nn}^{\rm ex}$. Neglecting this μ-dependence one gets the accustomed variation of the chemical potential

$$\delta\mu_n^0 = \delta U_1. \tag{14}$$

But taking of the second term into account yields a dramatic deviation from this traditional result when nuclei with a big neutron excess are considered.

Indeed, at small values of μ_n, the amplitude $f_{nn}^{\rm ex}$, Eq. (7), taken at $E = 2\mu_n$ is singular:

$$f_{nn}^{\rm ex} = \frac{\alpha}{\sqrt{E}} = \frac{\alpha}{\sqrt{2\mu_n}}. \tag{15}$$

Here we neglected the small, compared to μ_n, value of the virtual level energy of the T-matrix in the singlet channel. The variation of Eq. (15) yields:

$$\delta f_{nn}^{\rm ex} = -f_{nn}^{\rm ex}\frac{\delta\mu_n}{2\mu_n}. \tag{16}$$

Upon substituting Eqs. (12)–(14), and (16) into Eq. (11) one finds

$$\delta\mu_n = \frac{\delta\mu_n^0}{1 + V_0/(2\mu_n)}, \tag{17}$$

[3] This equation contains also an additional term provided the neutrons occupy a new j-level. All the consideration presented below remains valid in this case, too.

where a short notation $V_0 = \frac{1}{2}C_0\rho_0\frac{N^2}{A^2}f_{nn}^{\rm ex}$ for the first term of Eq. (10) is introduced.

It can be easily seen that the denominator of the relation (17) is noticeably greater than unit for nuclei with small value of μ_n. Let us, for example, calculate this quantity for two isotopes of tin in the vicinity of the old drip line ($A_{\max} = 176$ which is the common value to all calculations for the tin isotopes as far as we know). We use the values of μ_n which will be found below and the corresponding values of the amplitude $f_{nn}^{\rm ex}$ taken from Fig. 1. We take also the standard values of the normalization parameters: $C_0 = 300$ MeV $\cdot$ fm^3 and $\rho_0 = 0.16$ fm^{-3}. First, let us consider the ^{150}Sn nucleus. In this case, we have $f_{nn}^{\rm ex} = -1.4$, $V_0 = -15$ MeV, and $\mu_n = -3.4$ MeV. The substitution of these values into Eq. (17) yields $\delta\mu_n = \delta\mu_n^0/3.2$. The analogous calculation for the ^{200}Sn isotope ($\mu_n = -2.0$ MeV, $f_{nn}^{\rm ex} = -1.66$, $V_0 = -22.4$ MeV) results in $\delta\mu_n = \delta\mu_n^0/6.5$. We see that deviations from the traditional scheme are really large and are growing as soon as with the value of $|\mu_n|$ becomes less and less. This explains qualitatively why nuclei (e.g., ^{200}Sn) which are strongly unbound in traditional calculations could exist in our approach.

Going to the actual calculations incorporating the effect discussed above, we start from an advanced version of the SC FFS theory, the so-called quasiparticle Lagrange method (QLM).[4] It was devised in Ref. [8] for magic nuclei and extended in Refs. [15,16] to include superfluidity. All necessary modifications in the scheme can be explained for a simpler, nonsuperfluid case. We present here only a brief sketch of the QLM referring to Refs. [8,15,16] for details. The approach utilizes the Lagrange formalism which is the most convenient when the energy dependence effects are considered explicitly. The effective Lagrangian is constructed in such a way that its variation results in the Dyson equation,

$$\bigl(\varepsilon - \varepsilon_p^0 - \Sigma_q\bigr)G_q = 1, \tag{18}$$

for the quasiparticle Green function G_q with the quasiparticle mass operator

$$\Sigma_q(\mathbf{r},\mathbf{k},\varepsilon) = \Sigma_0(\mathbf{r}) + \frac{1}{(k_{\rm F}^0)^2}\mathbf{k}\Sigma_1(\mathbf{r})\mathbf{k} + \frac{\varepsilon}{\varepsilon_{\rm F}^0}\Sigma_2(\mathbf{r}), \tag{19}$$

[4] The term "quasiparticle" is used here in Landau's (not Bogolyubov's) sense.

where the normalization parameters are: $k_{\rm F}^0 = \pi/(mC_0)$, $\varepsilon_{\rm F}^0 = (k_{\rm F}^0)^2/(2m)$.

The first two terms of Σ_q are common to the HF theory with effective velocity dependent forces (e.g., Skyrme forces). The third term, $\Sigma_2(r)$, is due to energy dependence. It determines the coordinate dependent Z-factor:

$$Z(\mathbf{r}) = \bigl(1 - \Sigma_2(\mathbf{r})/\varepsilon_{\rm F}^0\bigr)^{-1}. \tag{20}$$

The solution of the Eq. (18) can be expressed in terms of the eigenfunctions Ψ_λ of the corresponding homogeneous equation as:

$$G_q(\mathrm{r}_1,\mathrm{r}_2,\varepsilon) = \sum_\lambda \frac{\Psi_\lambda^*(\mathrm{r}_1)\Psi_\lambda(\mathrm{r}_2)}{\varepsilon - \varepsilon_\lambda + i\delta\,\mathrm{sgn}(\varepsilon_\lambda - \mu)}, \tag{21}$$

where δ is small and positive. The functions Ψ_λ are orthonormalyzed with the weight:

$$\int d\mathbf{r}\,\Psi_\lambda^*(\mathbf{r})Z^{-1}(\mathbf{r})\Psi_{\lambda'}(\mathbf{r}) = \delta_{\lambda\lambda'}. \tag{22}$$

The quasiparticle density $\nu_0(\mathbf{r})$ associated with the Ψ_λ-functions is:

$$\nu_0(\mathbf{r}) = \sum_\lambda n_\lambda|\Psi_\lambda(\mathbf{r})|^2, \tag{23}$$

where $n_\lambda = (0,1)$ are the quasiparticle occupation numbers. It differs from the usual density $\rho(\mathbf{r})$ normalized to the total particle number by the Z-factor:

$$\nu_0(\mathbf{r}) = Z(\mathbf{r})\rho(\mathbf{r}). \tag{24}$$

There are two additional densities introduced in Ref. [8], the quasiparticle kinetic energy density

$$\nu_1(\mathbf{r}) = \frac{1}{(k_{\rm F}^0)^2}\sum_\lambda n_\lambda|\nabla\Psi_\lambda(\mathbf{r})|^2, \tag{25}$$

and the total quasiparticle energy density

$$\nu_2(\mathbf{r}) = \frac{1}{\varepsilon_{\rm F}^0}\sum_\lambda n_\lambda\varepsilon_\lambda|\Psi_\lambda(\mathbf{r})|^2, \tag{26}$$

The density $\nu_1(\mathbf{r})$ is analogous to the quantity $\tau(\mathbf{r})$ of the HF theory [7], whereas the density $\nu_2(\mathbf{r})$ is a new ingredient of the QLM, which does not appear in the HF approach.

In notation of Ref. [8], the density of the interaction Lagrangian L' is

$$\mathcal{L}'(\mathbf{r}) = -C_0\Bigg[\frac{\lambda_{00}}{2}\nu_{0+}^2(\mathbf{r}) + \frac{\lambda'_{00}}{2}\nu_{0-}^2(\mathbf{r}) + \frac{2\gamma}{3\rho_0^0}\nu_{0+}(\mathbf{r})\nu_{0n}(\mathbf{r})\nu_{0p}(\mathbf{r}) + \lambda_{01}\nu_{0+}(\mathbf{r})\nu_{1+}(\mathbf{r}) + \lambda'_{01}\nu_{0-}(\mathbf{r})\nu_{1-}(\mathbf{r}) + \lambda_{02}\nu_{0+}(\mathbf{r})\nu_{2+}(\mathbf{r}) - \frac{\lambda_{00}r_0^2}{2}\big(\nabla\nu_{0+}(\mathbf{r})\big)^2\Bigg] + \mathcal{L}_1, \qquad (27)$$

where the normalization density is $\rho_0^0 = 2(k_{\mathrm{F}}^0)^3/(3\pi^2)$, and $\mathcal{L}_1$ includes the Coulomb and the spin-dependent terms (mainly, the spin-orbit one). It is worth to mention that the isotopic structure of the "triple" term in Eq. (27) (proportional to γ) is similar to that of the Skyrme Hamiltonian of Ref. [7].

To extract from the Lagrangian L' the LM amplitude F, which, in accordance with the Landau prescription, is the second derivative of the ground state energy, one needs to find the corresponding Hamiltonian expressed in terms of the usual densities $\rho(\mathbf{r})$ and $\tau(\mathbf{r})$. Then it is necessary to calculate the second variational derivative with respect to ρ. This leads to a rather cumbersome expression which can be found in Ref. [13]. However, the relation of the constants $\lambda_{00}, \lambda'_{00}$ with the external LM amplitudes f^{ex}, f'^{ex} in Eqs. (3), (4) can be understood without using any explicit formula. Indeed, in the asymptotic region outside the nucleus the Z-factor is tending to unit and the densities $\nu_0(\mathbf{r})$ and $\rho(\mathbf{r})$ coincide. Therefore, the variation of L' with respect to ρ can be replaced by the variation with respect to ν_0. The amplitudes under consideration originate from the first three terms of Eq. (27). But the second variation of the third term vanishes at large r and one ends up with the following identities:

$$\lambda_{00} = f^{\mathrm{ex}}, \qquad \lambda'_{00} = f'^{\mathrm{ex}}. \qquad (28)$$

Of course, they follow also from the explicit relation for F of Ref. [13].

Dealing with superfluid nuclei, we utilize the modification of QLM developed for this case in Refs. [15, 16]. Although the method itself is, in principle, rather general, the practical scheme of Refs. [15,16] is quite simple. The main approximations are as follows. First, the λ-representation with a limited λ-basis is used ($\varepsilon_{\min} < \varepsilon_\lambda < \varepsilon_{\max}$, $\varepsilon_{\min} = -(20-25)$ MeV, $\varepsilon_{\max} = 5$ MeV), with the discretization of the continuum. Second, the δ-form density independent ("volume") pairing interaction is considered with the strength depending on the basis, in accordance with the prescription of Ref. [4]: $\Gamma_\xi = C_0 \ln^{-1}(C_p/\xi)$, where $\xi = \sqrt{(\mu - \varepsilon_{\min})(\varepsilon_{\max} - \mu)}$. At last, the diagonal approximation for the gap Δ is used: $\Delta_{\lambda\lambda'} = \Delta_\lambda \delta_{\lambda\lambda'}$. Thus, the method of describing the pairing effects in Refs. [15, 16] in the main points coincides with the BCS approximation. Although such a scheme possesses some well-known deficiencies for nuclei near the drip line and there exist much more advanced approaches for this case (see, e.g., Refs. [2,3]), we conserve here all details of the pairing scheme [15,16] (including the values of the pairing parameters) in order to separate the effect of the energy dependence of the mean field more clearly.

Our goal is to single out effects of the energy dependence of the external values of the LM amplitude. For convenience, let us rewrite the first three terms of Eq. (27) substituting the quantities f^{ex}, f'^{ex} for $\lambda_{00}, \lambda'_{00}$:

$$\mathcal{L}_0(\mathbf{r}) = -C_0\Bigg[\frac{f^{\mathrm{ex}}}{2}\,\nu_{0+}^2(\mathbf{r}) + \frac{f'^{\mathrm{ex}}}{2}\,\nu_{0-}^2(\mathbf{r}) + \frac{2\gamma}{3\rho_0^0}\nu_{0+}(\mathbf{r})', \nu_{0n}(\mathbf{r})\nu_{0p}(\mathbf{r})\Bigg]. \qquad (29)$$

If we consider f^{ex}, f'^{ex} as phenomenological parameters, there is no difference between this expression and the initial one, just the physical meaning of the constants $\lambda_{00}, \lambda'_{00}$ is more transparent. The next step is the calculation of these parameters from Eqs. (5), (6). They are $f^{\mathrm{ex}} = -2.6$ and $f'^{\mathrm{ex}} = 1.56$, Ref. [11], instead of $\lambda_{00} = -3.25$ and $\lambda'_{00} = 2.4$ of Ref. [8].[5] Now Eq. (29) contains only one adjustable parameter γ instead of the three those of the corresponding part of Eq. (27). We chose it in such a way to better reproduce the single-particle energy spectrum of the nucleus ^{124}Sn calculated in Ref. [15] where it was used for fitting the parameters of the calculation scheme and was found in a reasonable agreement with the experimental one. It should be noted that we do not analyze here the total binding energies and the density distributions but concentrate on the single-particle spectra

[5] These values of f^{ex}, f'^{ex} are closer to those of Ref. [3].

Table 1
Single-particle spectrum of the ^{124}Sn nucleus

λ	ε_λ, MeV		E_λ, MeV		
	[15]	This work	[15]	This work	Exp.
$2p_{1/2}$	−16.86	−17.94	−16.96	−18.05	
$1g_{9/2}$	−15.15	−15.93	−15.29	−16.11	
$2d_{5/2}$	−10.19	−10.82	−10.42	−11.05	
$3s_{1/2}$	−9.53	−8.82	−9.97	−9.18	−8.64
$1g_{7/2}$	−8.34	−8.32	−8.81	−9.19	−9.63
$2d_{3/2}$	−8.08	−8.05	−8.68	−8.66	−8.52
$1h_{11/2}$	−7.19	−6.87	−5.83	−5.08	−5.73
$2f_{7/2}$	−2.34	−2.29	−2.24	−2.14	
$3p_{3/2}$	−0.97	−0.96	−0.94	−0.91	
$3p_{1/2}$	−0.42	−0.26	−0.39	−0.23	

because they are mainly responsible for the position of the drip line. We found that the spectrum obtained with the ab initio constants $f^{\rm ex}$, $f'^{\rm ex}$ and with the value of $\gamma = 1.6$ (instead of $\gamma = 3.2$ in Refs. [8,15]) is in a reasonable agreement with that of Ref. [15] (see Table 1), where $E_\lambda = \mu \pm \sqrt{(\varepsilon_\lambda - \mu)^2 + \Delta_\lambda^2}$.

The generalization of Eq. (29) to nuclei with a large neutron excess, where the isotopic symmetry is violated, is quite obvious:

$$\begin{aligned}\tilde{\mathcal{L}}_0(\mathbf{r}) = C_0\Bigg[&\frac{1}{2}f_{nn}^{\rm ex}(E=2\mu_n)\nu_{0n}^2(\mathbf{r}) \\ &+ \frac{1}{2}f_{pp}^{\rm ex}(E=2\mu_p)\nu_{0p}^2(\mathbf{r}) \\ &+ f_{np}^{\rm ex}(E=\mu_n+\mu_p)\nu_{0n}(\mathbf{r})\nu_{0p}(\mathbf{r}) \\ &+ \frac{2\gamma}{3\rho_0^0}\nu_{0+}(\mathbf{r})\nu_{0n}(\mathbf{r})\nu_{0p}(\mathbf{r})\Bigg], \end{aligned} \tag{30}$$

where the external interaction LM amplitudes $f_{nn}^{\rm ex}$, $f_{pp}^{\rm ex}$, $f_{np}^{\rm ex}$ are energy dependent and should be calculated for the nucleus under consideration in the same way as in Eq. (9). Again the energy dependence of the first of them only is essential in the case under consideration. In the two next amplitudes, this dependence is retained just for presenting a more general form which is essential, e.g., for nuclei near the proton drip line.

We made a series of the self-consistent calculations for the chain of the tin isotopes, first, with the phenomenological Lagrangian, Eq. (27),[6] and, second, with the semi-microscopic one, Eq. (30). Results are displayed in Fig. 2, together with the predictions of

[6] They just repeat the calculations of Ref. [16].

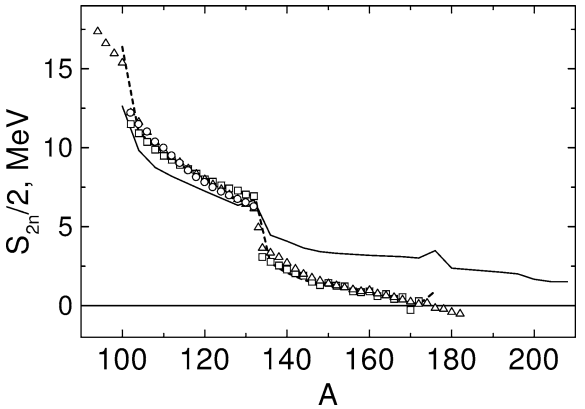

Fig. 2. A half of two-neutron separation energy $S_{2n}/2$ calculated for even–even tin isotopes. Results of the semi-microscopic calculations of this work (solid line) are displayed together with predictions of Ref. [16] (dashed line), Ref. [2] (triangles) and Ref. [3] (squares). Experimental values of $S_{2n}/2$ are shown by circles.

Dobaczewski et al. [2] and Fayans et al. [3]. The latter two approaches, just as the one of Ref. [16], use the phenomenological energy-independent forces.[7] The results of the three phenomenological calculations are quite close to each other. Qualitatively, they could be considered as one "phenomenological" curve which is in very good agreement with available experimental data. At small and intermediate values of the asymmetry parameter $y = (N - Z)/A$ the results of our semi-microscopic calculation deviate from the phenomenological curve (and, consequently, from the experimental data), but not significantly. The reason of this deficiency is quite simple. The matter is that the one-parameter Eq. (30) does not permit to obtain simultaneously reasonable values of the two "inner" amplitudes, $f^{\rm in}$ and $f'^{\rm in}$. It is well known that the Skyrme prescription of the isotopic structure of the "triple" term used in Eq. (30) is not obligatory for the effective force. The simplest two-parameter generalization is:

$$\begin{aligned}&\frac{2\gamma}{3\rho_0^0}\nu_{0+}(\mathbf{r})\nu_{0n}(\mathbf{r})\nu_{0p}(\mathbf{r}) \\ &\quad\to \frac{1}{6\rho_0^0}\big[\gamma\nu_{0+}^3(\mathbf{r}) + \gamma'\nu_{0+}(\mathbf{r})\nu_{0-}^2(\mathbf{r})\big], \end{aligned} \tag{31}$$

[7] The trivial, linear energy dependence, taken into account in Ref. [16], which is incorporated into Eq. (27) via the density $\nu_2(\mathbf{r})$, is not important for the effect under consideration.

which reduces to Eq. (30) at $\gamma' = -\gamma$. The ansatz of Eq. (31) seems to be more adequate to ideas underlying the approach suggested. Such a generalization, with a new adjustment of the parameters, will be carried out in a separate work. Preliminary estimates show that in this case the "semi-microscopic" results for stable nuclei become much closer to the phenomenological ones. On the other hand, the behaviour at large values of $N - Z$ is insensitive to possible modifications of the energy-independent parameters. Indeed, in terms of the qualitative consideration based on the SCR, such a variation results only in a change of the numerator of Eq. (17), whereas the denominator, which is responsible for the effect under consideration, remains unchanged. When we approach to the drip line, deviations grow, the semi-microscopic curve being significantly higher than the phenomenological one. It is remarkable that our calculation predicts the existence of nuclei beyond the commonly admitted end of the tin chain, $A_{\max} = 176$. It is worth to mention that this value of $A_{\max}$ corresponds to the value of $y = 0.43$ which exceeds the critical value $y_0 = 0.37$ predicted for the asymmetric nuclear matter [17]. The energy dependence effects under consideration make this difference between finite nuclei and nuclear matter more pronounced. We interrupted our calculations at $A = 208$ just because we are conscious that the deficiencies of the approach become more serious with the value of $|\mu_n|$ becoming less. One of them is too schematic consideration of pairing. Then, corrections to the standard FFS theory could appear when the energy dependence of the LM amplitude is significant.

It is worth to note that the two recipes of taking into account the energy dependence effects, Eqs. (9) and (30), are not identical. Indeed, in the first one the amplitudes f^{in}, f'^{in} are completely energy independent. As to the second recipe, the parameter γ is supposed to be energy independent that leads to some energy dependence of the quantities f^{in}, f'^{in} evaluated via the procedure described above. In the framework of a pure phenomenology, it is difficult to choose between the two ansatz. In our regular calculations, we use the second choice just because we have a well developed calculation scheme for this case. A more consistent microscopic theory for solving this problem is necessary. Such an approach based on the Brueckner theory for finite nuclei is now in progress [13,14], but it is yet far from to be completed. A semimicroscopic model for the scalar–isoscalar LM amplitude $f(\mathbf{r})$ suggested in Ref. [13] on the base of an approximate calculation of the G-matrix for a slab of nuclear matter [14] shows some energy dependence of the parameter f^{in}, but it is rather smooth. Evidently, the truth is somewhere between the two ansatz discussed above. What is important for our preliminary analysis, that both of them predict qualitatively the same. Namely, the energy dependence of the external LM amplitude f_{nn}^{ex} makes the neutron potential well deeper when the absolute value of μ_n becomes less. This effect could help nuclei with a huge neutron excess to survive.

The calculation scheme proposed in this Letter possesses some evident deficiencies partially discussed above. Improvement of them could make the effect less pronounced. That is the reason why the title of the Letter is formulated as a question. However, the effect itself is based on a transparent physical phenomenon. Therefore, we expect that it should remain, at least qualitatively, after all necessary corrections to the approximate calculation presented in this Letter.

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